

Re-born fireballs in Gamma-Ray Bursts

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ABSTRACT

We consider the interaction between a relativistic fireball and material assumed to be still located just outside the progenitor star. Only a small fraction of the expected mass is sufficient to efficiently decelerate the fireball, leading to dissipation of most of its kinetic energy. Since the scattering optical depths are still large at distances comparable to the progenitor radius, the dissipated energy is trapped in the system, accelerating it to relativistic velocities. The process resembles the birth of another fireball at radii $R \sim 10^{11}$ cm, not far from the transparency radius, and with a starting bulk Lorentz factors $\Gamma_c \sim 10$. As seen in the observer frame, this “re-generated” fireball appears collimated within an angle $\theta_j = 1/\Gamma_c$. If the central engine works intermittently, the funnel can, at least partially, refill and the process can repeat itself. We discuss how this idea can help solving some open issues of the more conventional internal shock scenario for interpreting the Gamma-Ray Burst properties.

Key words: gamma rays: bursts, X-rays: general, radiation mechanisms: general.

1 INTRODUCTION

The internal/external shock scenario (see e.g. Piran 2004; Meszaros 2006) is currently the leading model to explain the complex phenomenology of Gamma-Ray Bursts (GRBs) prompt and afterglow emission. Despite it can account for many observed characteristics, there are a few open issues and difficulties that this model cannot solve, or can accommodate only with some important modifications.

Here we recall some problems of the standard scenario and mention some ideas already put forward to account for them.

- *Efficiency I: high Γ -contrast* — In internal shocks only the relative kinetic energy of the two colliding shells can be dissipated. Thus “dynamical” efficiencies of only a few per cent can be achieved for colliding shells whose Lorentz factors Γ differ by a factor of order unity. Such efficiency has to include energy dissipated into randomizing protons, amplifying (or even generating) magnetic fields and accelerating emitting leptons. As the emitted radiation is produced only by the latter component, it corresponds to just a fraction of the dynamical efficiency. This problem, pointed out, among others, by Kumar (1999) can be solved by postulating contrasts in Γ much exceeding 100 (Belobodorov 2000; Kobayashi & Sari 2001). In these cases the typical Lorentz factor of GRBs should thus largely exceed the “canonical” value ~ 100 . In this case it is difficult to understand how the value of the peak energy of the prompt spectrum does not wildly change.

- *Efficiency II: Afterglow/Prompt power ratio* — A related inconsistency concerns the observed ratios of the bolometric fluence originated in the afterglow to that in the prompt phase. Since exter-

nal shocks are dynamically more efficient than internal ones, such ratio is expected to exceed one, contrary to current estimates. The problem has been exacerbated by the recent observations by the Swift satellite, showing that the X-ray afterglow light curve seen after a few hours – thought to be smoothly connected with the end of the prompt – comprises a steep early phase. As a consequence, the total afterglow energy is less than what postulated before. Willingale et al. (2007), parameterizing the behavior of the Swift GRB X-ray light curves, derived an average X-ray afterglow-to-prompt fluence ratio around 10 per cent (see also Zhang et al. 2007). Furthermore, the very same origin of the early X-ray radiation as produced by external shocks is questioned, since its behavior is different from the optical one (e.g. Panaitescu et al. 2006). If the X-ray emission does not originate in the afterglow phase, this further reduces the above ratio.

- *The spectral energy correlations* — Correlations have been found between i) the energy where most of the prompt power is emitted (E_{peak}) and the isotropic prompt bolometric energetics $E_{\gamma,\text{iso}}$ (Amati et al. 2002; Amati 2006), and ii) E_{peak} and the collimation corrected energetics E_γ . The slope of the former correlation is $E_{\text{peak}} \propto E_{\gamma,\text{iso}}^{1/2}$ while the slope of the latter depends on the radial profile of the circumburst density. For a homogeneous density medium $E_{\text{peak}} \propto E_\gamma^{0.7}$ (Ghirlanda et al. 2004), while for a wind-like profile in density ($\propto r^{-2}$) the correlation is linear: $E_{\text{peak}} \propto E_\gamma$ (Nava et al. 2006; Ghirlanda et al. 2007a). If linear, the relation is Lorentz invariant and indicates that different GRBs roughly emit the same number of photons at the peak (i.e. $E_\gamma/E_{\text{peak}} \sim \text{constant}$).

Note that the derivation of E_γ requires not only information on

the jet break time t_j , but also a model relating t_j with the collimation angle θ_j , which in turn depends on the circumburst density value, profile and the radiative efficiency η (i.e. $E_\gamma = \eta E_{\text{kin}}$). The phenomenological connection among the three observables $E_{\gamma, \text{iso}}$, E_{peak} and t_j , as found by Liang & Zhang (2005), is instead model-independent. It is of the form $E_{\text{iso}} \propto E_{\text{peak}}^a t_j^{-b}$, which for $b \sim -1$ is consistent with the Ghirlanda relation (both in the homogeneous and wind case; see Nava et al. 2006). A further tight phenomenological relation appears to link three prompt emission quantities: the isotropic peak luminosity L_{iso} , E_{peak} and the time interval $T_{0.45}$ during which the emission is above a certain level (Firmani et al. 2006). All these correlations were not predicted by the internal/external shock scenario, and can only be reconciled with it as long as specific dependences of the bulk Lorentz factor upon E_{iso} are satisfied (see Table 1 in Zhang & Meszaros 2002).

The above issues motivate the search for alternatives or for substantial modifications of the standard model. The efficiency problem and the existence of the spectral-energy relations prompted Thompson (2006, T06 hereafter) and Thompson Meszaros & Rees (2007, T07 hereafter) to suggest that, besides internal shocks, dissipation might also occur because of the interaction between the fireball and the walls of the funnel in the star through which it propagates (see §2). This hypothesis also introduces a typical scale to the problem, namely the radius of the progenitor star ($R_* \sim 10^{10}\text{--}10^{11}$ cm): shear instabilities within R_* can reconvert a significant fraction of bulk kinetic energy into heat. Since this dissipation occurs up to R_* (i.e. not far from the transparency radius), the increased internal energy can only partially reconvert into bulk motion via adiabatic expansion, increasing the efficiency. Similarly, studies of magnetized fireballs (e.g. Drenkhahn & Spruit, 2002; Giannios & Spruit 2007) have shown that dissipation of magnetic energy through reconnection can also contribute to increase the radiation content of the fireball at relatively large radii.

Along the above mentioned lines, in this Letter we propose a further possible way in which a large fraction of the fireball bulk energy can be dissipated at distances $R \sim R_*$. This assumes that at these distances the fireball collides with some mass which is (nearly) at rest: a small fraction of the mass swept up in the funnel left along the fireball propagation axis is sufficient to lead to efficient dissipation. Hereafter such mass will be referred to as “IDM” (Intervening Debris of the cocoon Material).

The treatment of the collision is simplified, in order to allow an analytical and simple description. We assume the IDM to be at rest and homogeneous in density and the fireball to have a Lorentz factor $\Gamma \gg 1$. The interaction is described in “steps”, while in reality will be continuous in time. A complete treatment of the dynamics and emission properties of our model requires numerical simulations (of the kind presented in Morsony et al. 2007, introducing some erratic behavior of the injected jet energy). Interestingly, the model predicts that jet properties depend on the polar angle (like in a structured jet, see Rossi et al. 2002) and it naturally implies a connection among the observed spectral-energy correlations.

2 SHEAR-DRIVEN INSTABILITIES AND DISSIPATION OF BULK KINETIC ENERGY

As mentioned, T06 and T07 proposed a model in which the efficiency of the dissipation of kinetic energy into radiation is enhanced with respect to the internal shock scenario and the spectral energy correlations, in particular the Amati one, can be accounted

for. At the same time, in their scenario synchrotron emission could play a minor role, the radiation field being dominated by thermalized high energy photons or by the inverse Compton process.

For what follows it is useful to summarize their main arguments here. Consider a fireball that at some distance $R_0 \lesssim R_*$ from the central engine is moving relativistically with a bulk Lorentz factor Γ_0 . The fireball is initially propagating inside the funnel of the progenitor star. T06 and T07 assume that a large fraction of the energy dissipated at R_0 is thermalized into black-body radiation of luminosity

$$L_{\text{BB,iso}} = 4\pi R_0^2 \Gamma_0^2 \sigma T_0'^4 = 4\pi \frac{R_0^2}{\Gamma_0^2} \sigma T_0^4, \quad (1)$$

where T_0' and $T_0 = \Gamma_0 T_0'$ are the temperatures at R_0 in the comoving and observing frame, respectively. The collimation corrected luminosity is $L_{\text{BB}} = (1 - \cos \theta_j) L_{\text{BB,iso}}$ which, for small semi-aperture angles θ_j of the jetted fireball (assumed conical), gives

$$\theta_j^2 \sim \frac{2L_{\text{BB}}}{L_{\text{BB,iso}}}. \quad (2)$$

A key assumption of the model is that $\Gamma_0 \sim 1/\theta_j$. The argument behind it is that if $\Gamma_0 \gg 1/\theta_j$, shear-driven instabilities have not time to grow (in the comoving frame), while in the opposite case the flow mixes easily with the heavier material and decelerates to $\Gamma \sim 1$. Then, assuming $\Gamma_0 = 1/\theta_j$ and substituting it in Eq. 1

$$L_{\text{BB,iso}} \sim 8\pi R_0^2 \frac{L_{\text{BB}}}{L_{\text{BB,iso}}} \sigma T_0^4. \quad (3)$$

Setting $E_{\text{BB,iso}} = L_{\text{BB,iso}} t_{\text{burst}}$ and $E_{\text{BB}} = L_{\text{BB}} t_{\text{burst}}$, where t_{burst} is the duration of the prompt emission, gives

$$E_{\text{peak}} \propto T_0 \propto E_{\text{BB,iso}}^{1/2} E_{\text{BB}}^{-1/4} t_{\text{burst}}^{-1/4}. \quad (4)$$

This corresponds to the Amati relation if E_{BB} is similar in different bursts and the dispersion of GRB durations is also limited. It should be noted that a relation similar to the Amati one can be also recovered adopting $E_{\text{BB}} \propto E_{\text{peak}}^a$, as suggested by the Ghirlanda relation. For instance, for $a = 1$ (wind case)

$$E_{\text{peak}} \propto E_{\text{BB,iso}}^{2/5} t_{\text{burst}}^{-1/5}. \quad (5)$$

For the derivation of Eq. 4 (and 5) a key assumption is the dependence on temperature of the black-body law, which leads to both a slope and a normalization similar to those characterizing the Amati relation.

We can ask what happens if, instead of a black-body, one assumes that the spectrum is a cutoff power-law. This question is particularly relevant since the burst spectrum is rarely described by a pure black-body (even if some bursts is, see Ghirlanda et al. 2003), and also the black-body plus power-law model (Ryde 2005) faces severe problems, even considering time resolved spectra (Ghirlanda et al. 2007b).

How is the above derivation modified if, instead of a black-body emission, the spectrum is best described by a cutoff power-law? Consider then a spectrum described in the comoving frame by $L'_{\gamma, \text{iso}}(E') \propto E'^{-\beta} \exp(-E'/E'_0)$ and approximate the observed isotropic bolometric luminosity as

$$L_{\text{iso}} \propto \Gamma_0^2 \left(\frac{E_{\text{peak}}}{\Gamma_0} \right)^{1-\beta} \propto \left(\frac{L_\gamma}{L_{\gamma, \text{iso}}} \right)^{-(1+\beta)/2} E_{\text{peak}}^{1-\beta}, \quad (6)$$

This leads to

$$E_{\text{peak}} \propto E_{\text{iso}}^{1/2} E_\gamma^{(1+\beta)/(2-2\beta)} t_{\text{burst}}^{-1/(1-\beta)}. \quad (7)$$

Therefore the dependence of the Amati relation can be recovered even for cutoff power-law spectra, but the normalization in this case is not determined. Note also that for $\beta = -3$ (Wien spectrum), Eq. 7 and Eq. 4 have the same dependences.

3 FIREBALL-IDM COLLISION

To excavate a funnel inside a progenitor star of mass $M_* = 10M_{*,1}$ solar masses, the “proto-jet” has to push out the mass that did not fall into the newly born black hole. This means a fraction of $(1 - \cos \theta_f)M_* = 0.1M_{*,1}\theta_{f,-1}^2$, where $\theta_f = 0.1\theta_{f,-1}$ is the funnel opening angle. This mass expands sideways as the proto-jet breaks out at the surface of the progenitor, forming a cocoon (see also Ramirez-Ruiz, Celotti & Rees 2002). We must expect, however, that after the break-out the region in front of the funnel will not be perfectly cleared of mass. To be irrelevant, the mass M_c left as IDM should be $\ll E_0/(\Gamma_0 c^2) \sim 5 \times 10^{-7} M_\odot$: less than one in a million particles should remain there.

If the jetted fireball is not continuous, this mass may be still there at the moment of the arrival of the new fireball pulse. And even if M_c is $(10^{-3}-10^{-4})$ the excavated mass, the IDM can have important dynamical effects on it. As the bulk velocity and energy content of the IDM can be neglected in comparison to those of the coming fireball, it will be approximated as initially at rest and cold.

An interesting aspect of this scenario concerns multi-peaked bursts, especially when pulses in the prompt emission are separated by quiescent periods. If the central engine works at a reduced rate during quiescence, material from the walls of the funnel, previously in pressure equilibrium with the jet, will tend to refill again the funnel (see also Wang & Meszaros 2007). This requires a time $T_f \sim \theta_f R/(\beta_s c) \sim 3 \times 10^{-4} (R/R_{0,7}) \theta_{f,-1}/\beta_{s,-1}$ s, where $\beta_s c$ is the sound speed, and $R_0 = 10^7 R_{0,7}$ cm is the radius at the base of the funnel. Then the amount of mass which should be pushed out again depends upon the quiescent time, but that after ~ 1 second, the funnel is completely closed, and the process repeats itself.

3.1 Results: dynamics and dissipation

In the following we derive the main characteristics of the system formed by a fireball impacting against the IDM. Consider a fireball with energy E_0 , mass M_0 and bulk Lorentz factor $\Gamma_0 = E_0/(M_0 c^2)$, impacting against a mass M_c , initially at rest.

The energy and momentum conservation laws read

$$\begin{aligned} M_0 \Gamma_0 + M_c &= \Gamma_c (M_0 + M_c + \epsilon'/c^2) \\ M_0 \Gamma_0 \beta_0 &= \Gamma_c \beta_c (M_0 + M_c + \epsilon'/c^2) \end{aligned} \quad (8)$$

where ϵ' is the dynamical efficiency measured in the frame moving at $\beta_c c$. Solving for β_c and ϵ' gives

$$\beta_c = \beta_0 \frac{M_0 \Gamma_0}{M_0 \Gamma_0 + M_c} \equiv \frac{\beta_0}{1+x}; \quad x \equiv \frac{M_c c^2}{M_0 \Gamma_0 c^2} \quad (9)$$

$$\epsilon' = \frac{E_0}{\Gamma_c} \left(1 + x - x \Gamma_c - \frac{\Gamma_c}{\Gamma_0} \right) \quad (10)$$

$$\epsilon \equiv \Gamma_c \epsilon'. \quad (11)$$

The Lorentz factor Γ_c is shown in Fig. 1 as a function of M_c for four values of E_0 (top panel) and as a function of E_0 for three values of M_c (bottom panel). One can see that for the process to be interesting (i.e. Γ_c significantly smaller than Γ_0) and to avoid “over-loading” of baryons (too small Γ_c) M_c is required to be in

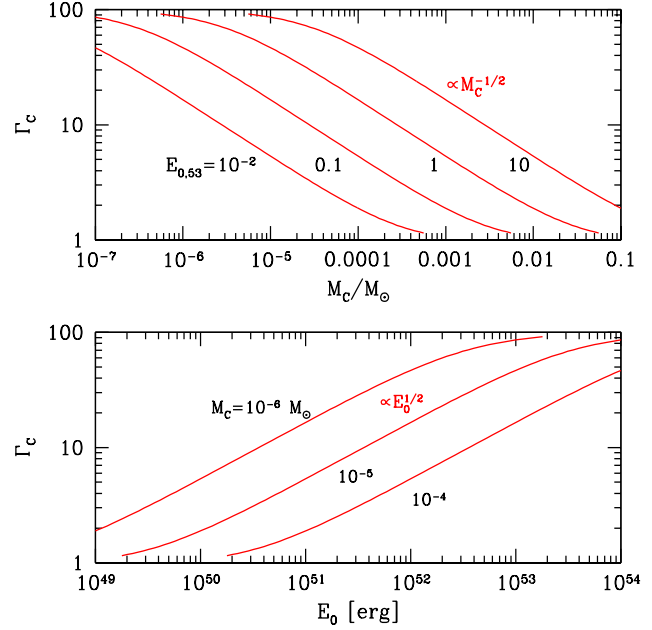


Figure 1. The Lorentz factor Γ_c of the IDM+fireball system as a function of M_c for different values of E_0 (top panel) and as a function of E_0 for selected values of M_c (bottom panel).

specific ranges, that depend on E_0 . These ranges, however, encompass almost two orders of magnitude.

Γ_c as a function of x is reported in Fig. 2. The power-law dependence $\Gamma_c \propto x^{-1/2}$ can be derived directly from Eq. 9, since $\Gamma_c^{-2} = 1 - \beta_c^2 \propto 2x$ for $x \ll 1$ and $\beta_0 \rightarrow 1$. Γ_c is limited to ~ 10 even for $x \sim 4 \times 10^{-3}$, corresponding to $M_c = 2.2 \times 10^{-5} E_{0,52} M_\odot$.

In Fig. 2 we show $\eta' \equiv \epsilon'/E_0$ and $\eta \equiv \epsilon/E_0$ as a function of x . For clarity also the fraction of the initial kinetic energy preserved after the collision

$$\eta_{\text{kin}} \equiv \frac{(\Gamma_c - 1)(M_0 + M_c)c^2}{E_0} = (\Gamma_c - 1) \left(x + \frac{1}{\Gamma_0} \right) \quad (12)$$

is plotted. The sum $\eta + \eta_{\text{kin}}$ is unity by definition.

Note that since x is a ratio, M_c and E_0 can be taken both as “isotropic” or “real” (collimation corrected) values.

3.2 Evolution of the fireball+IDM system

After the collision/dissipation phase the fireball+IDM is expected to be optically thick: the Thompson scattering optical depth of the IDM material is

$$\tau_c = \sigma_T n \Delta R = \frac{\sigma_T M_c}{4\pi R_c^2 m_p} = 3.2 \times 10^3 \frac{M_{c,-6}}{R_{c,11}^2}, \quad (13)$$

while the optical depth of the fireball just before the collision is of order

$$\tau_0 = \frac{\sigma_T E_0}{4\pi R^2 \Gamma_0 m_p c^2} = 1.8 \times 10^5 \frac{E_{F,52}}{R_{11}^2}, \quad (14)$$

where M_c and E_0 are here isotropic quantities. These large optical depths imply that the radiation produced following the collision is trapped inside the fireball+IDM system, which will expand because of the internal pressure. In the frame moving with Γ_c , the expansion is isotropic. In this frame some final Γ' will be reached. As seen in the observer frame, the expansion is highly asymmetric, and the

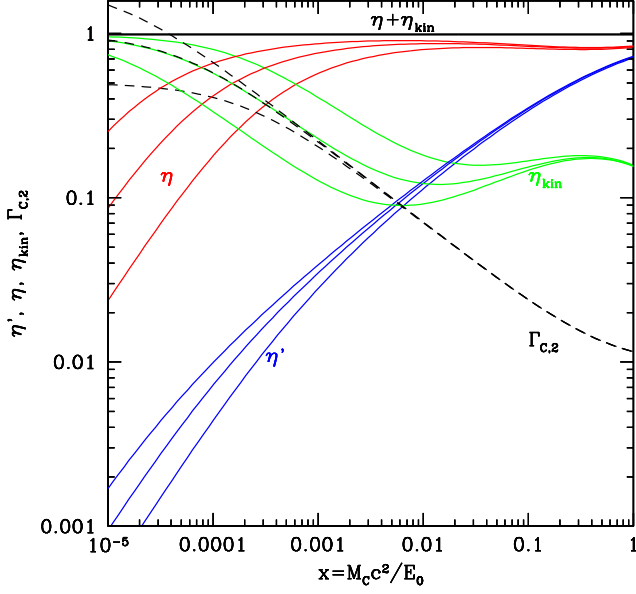


Figure 2. The Lorentz factor $\Gamma_{c,2} \equiv \Gamma_c/100$, the dynamical efficiency ϵ' and ϵ (in the comoving and observer frame, respectively) as a function of x . Each quantity is calculated for $\Gamma_0 = 50, 100$ and 200 (from top to bottom).

geometry of the system resembles a cone, with semi-aperture angle given by (Barbiellini, Celotti & Longo 2003)

$$\tan \theta_j = \frac{\beta_{\perp}}{\beta_{\parallel}} = \frac{\beta' \sin \theta'}{\Gamma_c (\beta' \cos \theta' + \beta_c)} \rightarrow \theta_j \sim \frac{1}{\Gamma_c}, \quad (15)$$

where the last equality assumed $\theta' = 90^\circ$ ($\beta' \sim 1$ and $\beta_c \sim 1$). Therefore the aperture angle of the re-born fireball is related to Γ_c , independently of the initial Γ -factor of the fireball *before* the collision or *after* the expansion. Note that the initial aperture angle of the fireball is irrelevant, as it is the aperture angle of the funnel of the progenitor star.

The fireball is not collimated in a perfect cone, and mass and energy propagate also outside θ_j . Since in the frame moving with Γ_c it expands isotropically, $M'(\Omega') = M'/(4\pi)$ is approximately constant. Therefore in the observer frame

$$M(\theta) = \frac{M'}{4\pi} \frac{d \cos \theta'}{d \cos \theta}. \quad (16)$$

Also the resulting Lorentz factor Γ is angle dependent: from the relativistic composition of velocity (e.g. Rybicki & Lightman 1979):

$$\begin{aligned} \beta_{\parallel} &= \beta \cos \theta = \frac{\beta' \cos \theta' + \beta_c}{(1 + \beta_c \beta' \cos \theta')} \\ \beta_{\perp} &= \beta \sin \theta = \frac{\beta' \sin \theta'}{\Gamma_c (1 + \beta_c \beta' \cos \theta')} \\ \beta(\theta) &= (\beta_{\parallel}^2 + \beta_{\perp}^2)^{1/2} \\ \Gamma(\theta) &= [1 - \beta^2(\theta)]^{-1/2} = \Gamma' \Gamma_c (1 + \beta' \beta_c \cos \theta'), \end{aligned} \quad (17)$$

The observed Γ -factor is constant up to angles slightly smaller than $1/\Gamma_c$, and decreases as θ^{-2} above.

As a consequence of the angular dependence of mass and bulk Lorentz factor, also the energy depends on θ as

$$E(\theta) = \Gamma(\theta) M(\theta) c^2. \quad (18)$$

Such dependences of mass, Γ and E on polar angle are illustrated in Fig. 3. The jet is structured and well approximated by a top-hat

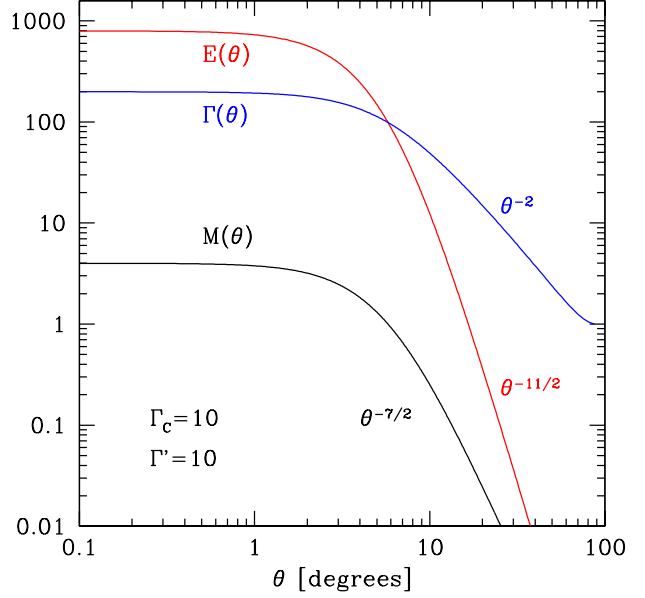


Figure 3. The mass, energy and bulk Lorentz factor (observer frame) as functions of the angle from the jet axis θ , for $\Gamma_c = 10$ and $\Gamma' = 10$. All the three quantities are constant up to $\theta \simeq 1/\Gamma_c$ and then decrease (approximately) as power-laws.

jet: the energy profile is nearly constant within an angle slightly smaller than $1/\Gamma_c$, and at larger angles decreases approximately as a steep power-law $E(\theta) \propto \theta^{-11/2}$, since $M(\theta) \propto \theta^{-7/2}$ and $\Gamma(\theta) \propto \theta^{-2}$ for $\theta \gg 1/\Gamma_c$. This particular behavior gives raise to an afterglow light curve indistinguishable from a top-hat jet (see Rossi et al. 2004).

4 DISCUSSION AND CONCLUSIONS

The most appealing feature of the proposed model is the high efficiency in re-converting the fireball kinetic energy into internal pressure at a radius comparable to the radius of the progenitor star, i.e. on a scale not far from the transparency one. Also the observed energetics of the internal radiation will be large, since the system becomes transparent during (or slightly after) the expansion/acceleration phase, similarly to the standard fireball models of initially high entropy, where the fireball becomes transparent before coasting. Our model therefore increases the parameter space of high efficiency regimes. Our model is also similar to the model proposed by T06 and T07 and to those in which the dissipation of an energetically important magnetic field occurs at large radii (see e.g. Giannios & Spruit 2007). The efficiency of the energy re-conversion for the fireball-IDM collision is of the order of 50–80 per cent for large ranges in the mass of the IDM and energy of the fireball (see Fig. 2).

With respect to the idea proposed by T07, summarised in §2, in our scenario there is no requirement on any specific value for the fireball bulk Lorentz factor Γ_0 prior to its collision with the IDM. Note also that in the T07 model, a “standard” fireball (i.e. not magnetic) moving with $\Gamma_0 \sim 1/\theta_j$ can dissipate part of its kinetic energy, but it cannot reach final Γ -factors larger than Γ_0 , which is bound to be small for typical θ_j . In our case the final Γ can instead be large, (even if always smaller than Γ_0). The jet angle is $\sim 1/\Gamma_c$

even for large values of the initial Γ_0 and final Γ . This is a result of our model, and not an assumption.

The re-born fireball is structured and the $M(\theta)$ and $\Gamma(\theta)$ behaviors imply that $E(\theta)$ depends on θ as steep power-laws. Despite the angle dependence of the energy, the jet should produce an afterglow indistinguishable from a top-hat jet. Clearly our description of the fireball-IDM interaction is extremely simplified, aimed at building a physical intuition based on the analytical treatment. More realistic situations should be studied via numerical simulations, but we would like to comment on two aspects. i) Even if the IDM is initially at rest, as soon as the fireball starts depositing a fraction of energy and momentum, the IDM would begin to move. The whole process will take long enough that towards the end it would be probably better described as the interaction with a moving IDM, with a consequent loss of efficiency. ii) In a ‘continuous’ (non-intermittent) scenario, the IDM would predominantly interact with the fireball edge, causing only a partial dissipation of its energy and the formation of a “fast spine-slow layer” structure. In this case the determination of the relevant jet opening angle (i.e. within which most of the energy is concentrated) requires a more accurate numerical treatment.

Despite of these caveats, it is still interesting to consider whether the model can account for the spectral-energy correlations, or at least highlight the relations between them. Consider the Ghirlanda correlation in the wind case, $E_{\text{peak}} \propto E_\gamma$, that can be rewritten as $E_{\text{peak}} \propto \theta_j^2 E_{\gamma, \text{iso}}$ for small θ_j . The requirement that also $E_{\text{peak}} \propto E_{\gamma, \text{iso}}^{1/2}$ (Amati relation) leads to:

$$E_\gamma \theta_j^2 = \text{constant}, \quad (19)$$

i.e. more energetic bursts are more collimated, as predicted in our model. The above condition (Eq. 19) can be quantitatively satisfied if i) the mass against which the fireball collides is similar in different GRBs, namely $\Gamma_c \sim 1/\theta_j \sim E_0^{1/2}$ (see Fig. 1) and ii) the prompt emission luminosity E_γ is also a constant fraction of E_0 for different GRBs. Eq. 19 does not explain the Amati or Ghirlanda relations, although it offers some physical meanings of required connection between the two. Note that, in the standard internal shock model, one can recover the Amati relation if $\Gamma \sim \text{constant}$ (see Zhang & Meszaros 2002).

In this Letter we did not discuss the characteristics of the spectrum predicted in our scenario (Nava et al., in prep). In general terms, the most effective radiation process would be “dynamical Compton”, or Fermi “acceleration” of photons, as discussed by Gruzinov & Meszaros (2000) in the context of internal shocks. The large number of photons per proton in the fireball (corresponding to the “fossil” radiation that has accelerated the fireball itself in the first place) implies that a significant fraction of the internal energy following the fireball-IDM collision *directly* energizes photons, i.e. photons amplify their energy by interacting with leptons with bulk momentum not yet randomized in the shock. The process is analogous to particle acceleration in shocks and gives raise to high energy photons, conserving their number.

Since this occurs at $R \sim 10^{11}$ cm, i.e. slightly above the progenitor star, the re-born fireball will become transparent during or just after the acceleration phase (at a transparency radius $R_\tau \sim 3 \times 10^{12}$ cm, Eq. 13). This ensures that photons do not have time to loose energy via adiabatic expansion, once again leading to large efficiencies.

The timescale for the refilling of the funnel during quiescent phases of the central engine can be short enough that the process repeats itself. The estimated refilling timescale is of the order of

a second: bursts with shorter or longer ‘quiescent’ phases should then show different properties. If the fireball-IDM collision is the dominant process for the dissipation (and emission), more energetic spikes are expected to follow longer quiescent phases (see Ramirez-Ruiz & Merloni, 2001).

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